

FTUV/96-19

IFIC/96-21

# Nucleosynthesis constraints on heavy $\nu_\tau$ in the presence of annihilations to majorons

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*Contribution to the IV International Workshop  
on Theoretical and Phenomenological Aspects of  
Underground Physics, Toledo, Spain, 17-21 September 1995.*

## Abstract

We show that in the presence of sufficiently strong  $\nu_\tau$  annihilations to majorons, primordial nucleosynthesis constraints can not rule out any values of the  $\nu_\tau$  mass up the present laboratory limit.

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<sup>1</sup>Supported by Conselleria d'Educació i Ciència of Generalitat Valenciana. Work done in collaboration with J.C. Romão (IST, Lisbon), A.D. Dolgov and J.W.F. Valle (IFIC, València).

The tau-neutrino is the only one which can have mass in the MeV range. Among all studied limits on  $m_{\nu_\tau}$ , the kinematical ones are the only model-independent bounds. The last result from ALEPH [1] is

$$m_{\nu_\tau} < 23 \text{ MeV} \quad (1)$$

In addition to this, there are stronger bounds from cosmological considerations. From the contribution of stable  $\nu_\tau$  to the present relic density one gets  $m_{\nu_\tau} < 92\Omega h^2 \text{ eV}$  [2], where  $h = H_0/(100 \text{ kms}^{-1}\text{Mpc}^{-1})$ . This means that a massive  $\nu_\tau$  with mass in the MeV range must be unstable with lifetimes smaller than the age of the Universe.

Moreover, if massive  $\nu_\tau$ 's are stable during nucleosynthesis ( $\nu_\tau$  lifetime longer than  $\sim 100 \text{ sec}$ ), one can constrain their contribution to the total energy density, using the observed amount of primordial helium. This bound can be expressed through an effective number of massless neutrino species  $N_\nu$ . Using  $N_\nu < 3.4 - 3.6$ , the  $\nu_\tau$  mass range

$$0.5 \text{ MeV} < m_{\nu_\tau} < 35 \text{ MeV} \quad (2)$$

has been excluded [3, 4]. This forbids all  $\nu_\tau$  masses on the few MeV range. However, such masses are theoretically viable [5], and interesting for a possible solution of the structure formation problem as described in [6].

It is possible to weaken the constraints on  $\nu_\tau$  masses of (2), by adding new interactions of  $\nu_\tau$  beyond the standard ones. One may consider either that the  $\nu_\tau$ 's are unstable during nucleosynthesis [7] or that they possess new channels of annihilation beyond the standard ones. The last case is what we have considered.

We studied the effect of the presence of Majorana  $\nu_\tau$  annihilations to majorons (J). The interaction term of the Lagrangian that describes the coupling between neutrinos and majorons can be expressed [5] as  $\mathcal{L}_{\nu\nu\mathcal{J}} \sim \{\mathcal{J}\nu^\mathcal{T}\nu$ , where  $g$  is a  $3 \times 3$  matrix containing all possible couplings. In particular for  $\nu_\tau$  the main processes are those described in figure (1). We assume that the coupling constants  $g_{23}$  (non-diagonal) and  $g_{33}$  (diagonal) are such that during nucleosynthesis time only annihilations are important. In previous works only the decay processes were taken into account, though in general  $g_{23}$  and  $g_{33}$  may be simultaneously important.

We consider the interactions of massive  $\nu_\tau$ 's stable during the epoch just before nucleosynthesis. The  $\nu_\tau$ 's interact with leptons via the standard weak interactions,  $\nu_\tau \bar{\nu}_\tau \leftrightarrow \nu_0 \bar{\nu}_0$ ,  $e^+ e^-$ , as in [3, 4]. We assume that, in addition,

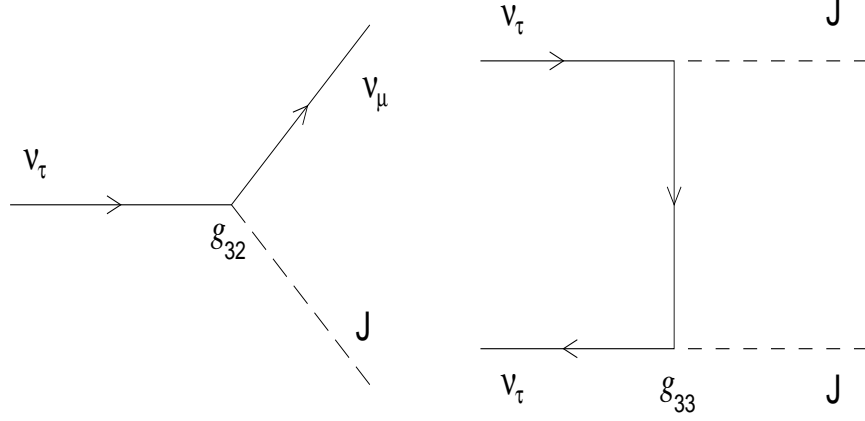


Figure 1: Non-diagonal and diagonal couplings of Majorana tau neutrinos and majorons.

the  $\nu_\tau$ 's annihilate to majorons via the diagonal coupling

$$\mathcal{L} \sim \} \mathcal{J} \nu_\tau^T \sigma_\tau \nu_\tau \quad (3)$$

where  $g \equiv g_{33}$  and  $\nu_\tau$  represents a two-component Majorana spinor. The corresponding elastic processes do not change the particle densities, but as long as they are effective they maintain all species with the same temperature.

Even after the  $\nu_\tau$ 's have decoupled from standard weak interactions, they remain in contact with majorons. Therefore, as happens with photons and neutrinos when the  $e^+e^-$  pairs annihilate, after the decoupling there are two plasmas with different temperatures, one constituted of  $\nu_\tau$ 's and  $J$ 's and the other of the rest of the particles.

The evolution of the number densities of tau neutrinos ( $n_\tau$ ) and majorons ( $n_J$ ) is given by the solution of a set of Boltzmann differential equations

$$\begin{aligned} \dot{n}_\tau + 3Hn_\tau = & - \sum_{i=e,\nu_0} \langle \sigma_i v \rangle (n_\tau^2 - (n_\tau^{eq})^2) \\ & - \langle \sigma_J v \rangle (n_\tau^2 - (n_\tau^{eq})^2 \frac{n_J^2}{(n_J^{eq})^2}) \end{aligned} \quad (4)$$

$$\dot{n}_J + 3Hn_J = \langle \sigma_J v \rangle (n_\tau^2 - (n_\tau^{eq})^2 \frac{n_J^2}{(n_J^{eq})^2}) \quad (5)$$

where  $\langle\sigma_\alpha v\rangle$  is the thermally averaged cross section of the annihilation process of  $\nu_\tau$ 's to particle-type  $\alpha$  [8], where  $\alpha = e, \nu_0, J$ . In the above equations we assumed Boltzmann statistics and  $n_i = n_i^{eq}$  for  $i = e, \nu_0$ .

Now let us briefly describe the calculations. First we normalized the number densities to the number density of massless neutrinos,  $n_0 \simeq 0.18T^3$ . We introduced  $r_\alpha \equiv n_\alpha/n_0$ , where  $\alpha = \nu_\tau, J$ , and the corresponding equilibrium functions  $r_\alpha^{eq}$ . On the other hand we performed the integrations using the dimension-less variable  $x \equiv m_{\nu_\tau}/T$ . The set of equations (4) and (5) is completed with the evolution of the  $\nu_\tau$  temperature, which is obtained from one of Einstein's equations

$$\dot{\rho} = -3H(\rho + P) . \quad (6)$$

Here  $\rho$  and  $P$  are the energy density and pressure of tau neutrinos and majorons, respectively.

We have integrated numerically the coupled differential equations in (4-6), obtaining the solutions of  $r_\tau$  for each pair of values  $(m_{\nu_\tau}, g)$ . The initial conditions are, for sufficiently high temperatures,  $r_\tau = r_\tau^{eq}$ ,  $r_J = r_J^{eq}$  and  $T_\tau = T_J = T_{\nu_0}$ .

The value of  $r_\tau(m_{\nu_\tau}, g)$  is used to estimate the variation of total energy density  $\rho_{tot} = \rho_R + \rho_{\nu_\tau}$ . In  $\rho_R$  all relativistic species are taken into account, including majorons and two massless neutrinos, whereas  $\rho_{\nu_\tau}$  is the energy density of massive  $\nu_\tau$ 's.

We can calculate now the effective number of massless neutrino species ( $N_\nu$ ) corresponding to each  $r_\tau(m_{\nu_\tau}, g)$ . First we run a program which calculates the evolution of the neutron fraction ( $r_n$ ), as presented e.g. in [9], varying the value of  $N_\nu$ . Then we incorporate  $\rho_{tot}$  to the same program and perform the integration for each pair of values  $(m_{\nu_\tau}, g)$ . Comparing  $r_n$  obtained in each case at  $T_\gamma \simeq 0.065 \text{ MeV}$  (the moment when practically all neutrons are wound up in  $^4\text{He}$ ), we can relate  $(m_{\nu_\tau}, g)$  and  $N_\nu$ .

The results are shown in figure (2). One can see that for a fixed  $N_\nu^{max}$ , a wide range of tau neutrino masses is allowed for large enough coupling constants  $g$ . One sees that all masses below 23 MeV are allowed by nucleosynthesis arguments, provided that the coupling between  $\nu_\tau$ 's and  $J$ 's exceeds a value of a few times  $10^{-4}$ . Such values are reasonable for many majoron models [5]. A more detailed study of this problem will be presented elsewhere [10].

Our conclusion is that the constraints on the mass of a Majorana  $\nu_\tau$  from primordial nucleosynthesis can be relaxed if annihilations  $\nu_\tau \bar{\nu}_\tau \leftrightarrow JJ$  are present.

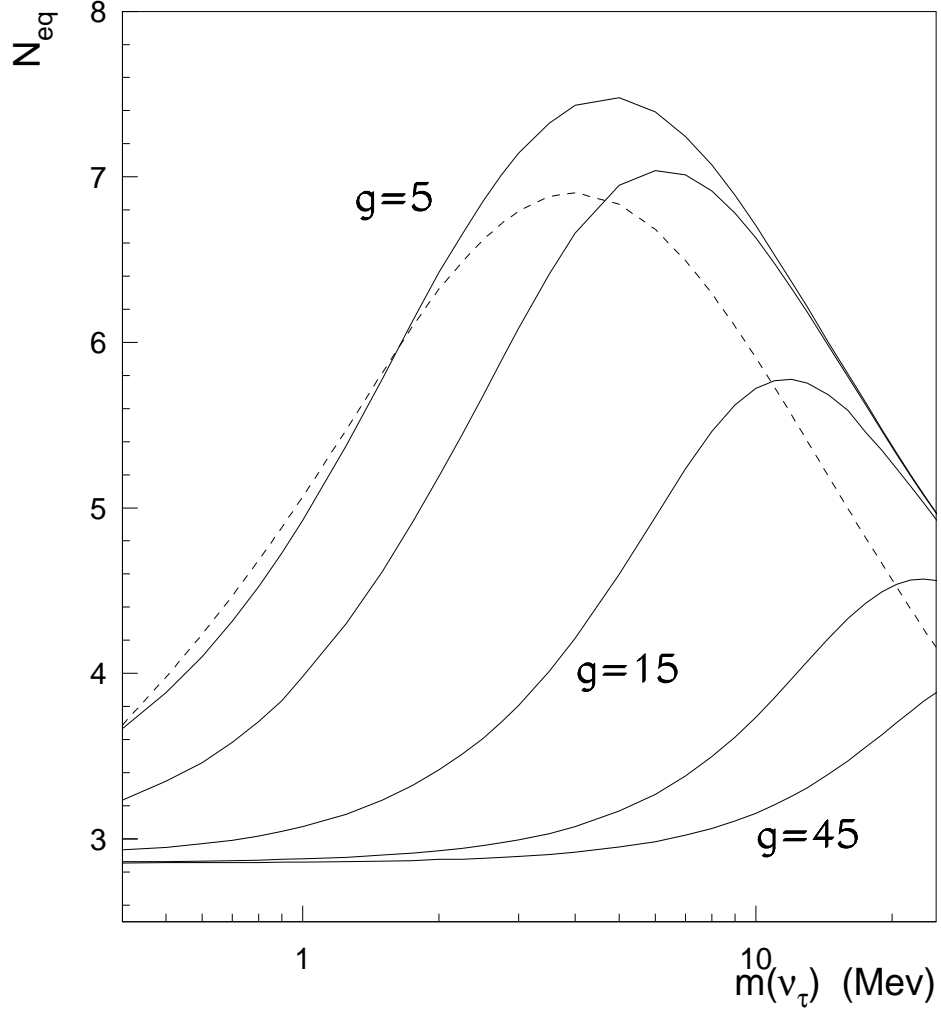


Figure 2: Effective number of massless neutrinos equivalent to the contribution of massive neutrinos with different values of  $g$  expressed in units of  $10^{-5}$ . For comparison, the dashed line corresponds to the case when  $g = 0$  and no majorons are present.

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